

# Unparticle effects in neutrino telescopes

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Recently H. Georgi has introduced the concept of unparticles in order to describe the low energy physics of a nontrivial scale invariant sector of an effective theory. We investigate its physical effects on the neutrino flux to be detected in a kilometer cubic neutrino telescope such as IceCube. We study the effects, on different observables, of the survival neutrino flux after through the Earth and the regeneration originated in the neutral currents. We calculate the contribution of unparticle physics to the neutrino-nucleon interaction and, then, to the observables in order to evaluate detectable effects in IceCube. Our results are compared with the bounds obtained by other non-underground experiments. Finally, the results are presented as an exclusion plot in the relevant parameters of the new physics stuff.

PACS numbers: PACS: 13.15.+g, 95.55.Vj

## I. INTRODUCTION

The Standard Model (SM) for the elementary particles interactions has been successfully tested at the level of quantum corrections. In particular high precision and collider experiments have tested the model and have placed the border line for physics effects at energies of the order of 1TeV [1]. On the other hand, new physics effect in the neutrino sector have recently received an important amount of experimental information coming from flavor oscillation [2]. This fact is the first evidence of neutrino masses different from zero, and hence, of physics beyond the SM. In this way, the neutrino sector and in particular neutrino-nucleon interactions, could be the place where new physics may become manifest again. Although the SM has been successful to describe the world at short distances, as a low energy effective theory of phenomena at higher scales, it leaves several open questions, e.g.: it does not predict the number of families and the fermions masses, has several free parameters, the mass generation mechanism through the Higgs boson, where its mass is not predicted, is untested and still leaves open the hierarchy problem. In these conditions, it is believed that we should have some kind of physics beyond the SM, which is called New Physics (NP)[1]. The search of NP proceeds mainly through the comparison of data with the SM predictions. The experimental way to look for NP effects in a model independent fashion is to construct observables that can be affected by this new physics and then compare the measurements with the mentioned SM expectation. Certain types of NP can already be present at the TeV scale and could participate in neutrino-nucleon interactions. Hence, these NP effects could possibly become apparent in neutrino telescopes.

In the other hand it is well-known that scale invariance has been a powerful tool in several branches of physics and the possibility of a scale invariant weak interacting sector with the low energy particle spectrum is not ruled out. In particular, H. Georgi [3] has proposed that a scale invariant sector which does not imply conformal invariance [4] with a non-trivial *IR* fixed point and coupled to the *SM* fields through the exchange of particles with a high mass scale  $M_U$  may appear much above the TeV energy scale. Below this energy scale, this sector induces unparticle operators  $\mathcal{O}_U$  with a non-integral scale dimension  $d_u$  that in turn have a mass spectrum which looks like a  $d_u$  number of massless

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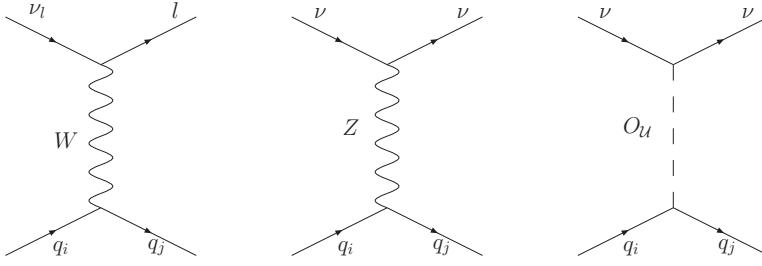


FIG. 1: Diagrams contributing to the neutrino-nucleon cross section.

particles. The couplings of these unparticles to the SM fields and in particular to standard neutrinos (i.e. massless and left-handed particles) and quarks are described by the effective Lagrangian [3]

$$\mathcal{L}_{eff} = \frac{\Lambda^{k+1-d_u}}{M_{\mathcal{U}}^k} \bar{f} \gamma_\mu (C_V + C_A \gamma_5) f \mathcal{O}_{\mathcal{U}}^\mu, \quad (1)$$

where  $\Lambda$  is the energy scale at which scale invariance emerges, the dimensionless coefficients  $C_{V(A)}$  are of order 1 for neutrinos and quarks and  $f$  is the generic fermion spinor. The operators with lowest possible dimension have the most important effect in the low energy effective theory regime. In Eq.(1) we have only included the vector unparticle operators  $\mathcal{O}_{\mathcal{U}}^\mu$  that couple with left neutrinos and both left and right quarks. Note that the left couplings to the neutrinos and quarks are taken equal.

In the present work we follow the Georgi's approach where the Feynman propagators of the unparticle operator  $\mathcal{O}_u^\mu$  is determined by the scale invariance [3],

$$\int d^4x e^{ipx} \langle 0 | T(O_U^\mu(x) O_U^\nu(0)) | 0 \rangle = i \frac{A_{d_u}}{2 \sin(d_u \pi)} \frac{-g^{\mu\nu} + p^\mu p^\nu / p^2}{(-p^2 - i\epsilon)^{2-d_u}}, \quad (2)$$

with

$$A_{d_u} = \frac{16 \pi^{5/2}}{(2\pi)^{2d_u}} \frac{\Gamma(d_u + \frac{1}{2})}{\Gamma(d_u - 1) \Gamma(2d_u)}. \quad (3)$$

The scale dimension  $d_u$  is restricted in the range  $1 < d_u < 2$ . Here, the condition  $d_u > 1$  is due to the non-integrable singularities in the decay rate [5] while  $d_u < 2$  is due to the convergence of the integrals [6].

## II. THE CROSS SECTION NEUTRINO-NUCLEON AND THE UNPARTICLES CONTRIBUTION.

In this section we consider the effective operator given in Eq.(1) and calculate their contribution to the neutrino-nucleon inclusive cross section:

$$\nu N \rightarrow \mu + \text{anything}, \quad (4)$$

where  $N \equiv \frac{n+p}{2}$  is an isoscalar nucleon. The corresponding process is pictured in Fig. 1 which has the SM charged and neutral current diagrams and the unparticle contribution. For charged currents the calculation is standard and we use it to compare our results with [7]. For neutral currents we have included the contributions of unparticles. The corresponding coupling and propagator are given in equations (1) and (2) respectively.

The SM results for the scattering amplitude for charged currents muon-neutrino scattering is

$$\mathcal{M}_{SM}^{CC} = -\frac{ig^2}{2(Q^2 + M_W^2)} \bar{l} \gamma_\mu P_L \nu \sum_{i=D,\bar{U}} \bar{q}_j \gamma^\mu P_L q_i, \quad (5)$$

and the corresponding differential cross section reads

$$\frac{d\sigma^{CC}}{dx dy} = \frac{G_F^2 s}{\pi} \left( \frac{M_W^2}{(Q^2 + M_W^2)} \right)^2 x [Q^{CC} + (1-y)^2 \bar{Q}^{CC}], \quad (6)$$

where for an isoscalar target we have the quark distribution functions

$$\begin{aligned} Q^{CC}(x, Q^2) &= \frac{u_v(x, Q^2) + d_v(x, Q^2)}{2} + \frac{u_s(x, Q^2) + d_s(x, Q^2)}{2} \\ &\quad + s_s(x, Q^2) + b_s(x, Q^2), \\ \bar{Q}^{CC}(x, Q^2) &= \frac{u_s(x, Q^2) + d_s(x, Q^2)}{2} + c_s(x, Q^2) + t_s(x, Q^2). \end{aligned} \tag{7}$$

Similarly, the SM neutral current amplitude is

$$\mathcal{M}_{SM}^{NC} = -\frac{ig^2}{2c_W^2(Q^2 + M_Z^2)} \bar{\nu} \gamma_\mu P_L \nu \sum_{i=U,D} \bar{q}_i \gamma^\mu (g_L^i P_L + g_R^i P_R) q_i, \tag{8}$$

where  $c_W = \cos \theta_W$ ,  $x_W = \sin^2 \theta_W$ ,  $g_L^U = 1/2 - 2x_W/3$ ,  $g_L^D = -1/2 + x_W/3$ ,  $g_R^U = -2x_W/3$  and  $g_R^D = x_W/3$ .

From the effective interaction Eq.(1) between unparticle and the SM fields we obtain the following four fermion amplitude

$$\mathcal{M}_{unp} = \frac{i}{Q^2} \frac{A_{du}}{2 \sin(d\pi)} \left( \frac{Q^2}{M_Z^2} \right)^{d_u-1} \bar{\nu} \gamma_\mu P_L \nu \sum_{i=U,D} \bar{q}_i \gamma^\mu (c_L^2 P_L + c_L c_R P_R) q_i, \tag{9}$$

where the left ( $c_L$ ) and right ( $c_R$ ) coupling constants are expressed in terms of the vector and axial vector coupling constants:

$$\begin{aligned} c_R &= c_V + c_A \\ c_L &= c_V - c_A, \end{aligned} \tag{10}$$

where

$$c_{V(A)} = \frac{C_{V(A)} \Lambda_u^{k+1-d_u}}{M_u^k M_Z^{1-d_u}}. \tag{11}$$

The total neutral current contribution, including the unparticle contribution, can be written as

$$\mathcal{M}^{NC} = -\frac{ig^2}{2c_W^2(Q^2 + M_Z^2)} \bar{\nu} \gamma_\mu P_L \nu \sum_{i=U,D} \bar{q}_i \gamma^\mu (\tilde{g}_L^i P_L + \tilde{g}_R^i P_R) q_i, \tag{12}$$

where

$$\begin{aligned} \tilde{g}_L^i &= g_L^i - (\delta g) c_L^2 \\ \tilde{g}_R^i &= g_R^i - (\delta g) c_L c_R, \end{aligned} \tag{13}$$

and

$$\delta g = \frac{A_{du}}{\sin(d_u \pi)} \frac{c_w^2}{g^2} \left( 1 + \frac{Q^2}{M_Z^2} \right) \left( \frac{Q^2}{M_Z^2} \right)^{(d_u-2)}. \tag{14}$$

The neutral current differential cross section is then

$$\begin{aligned} \frac{d\sigma^{NC}}{dxdy} &= \frac{G_F^2 s}{\pi} \left( \frac{M_Z^2}{Q^2 + M_Z^2} \right)^2 \sum_{i=U,D} x [\tilde{g}_L^{i2} (Q^i + (1-y)^2 \bar{Q}^i) \\ &\quad + \tilde{g}_R^{i2} (\bar{Q}^i + (1-y)^2 Q^i)], \end{aligned} \tag{15}$$

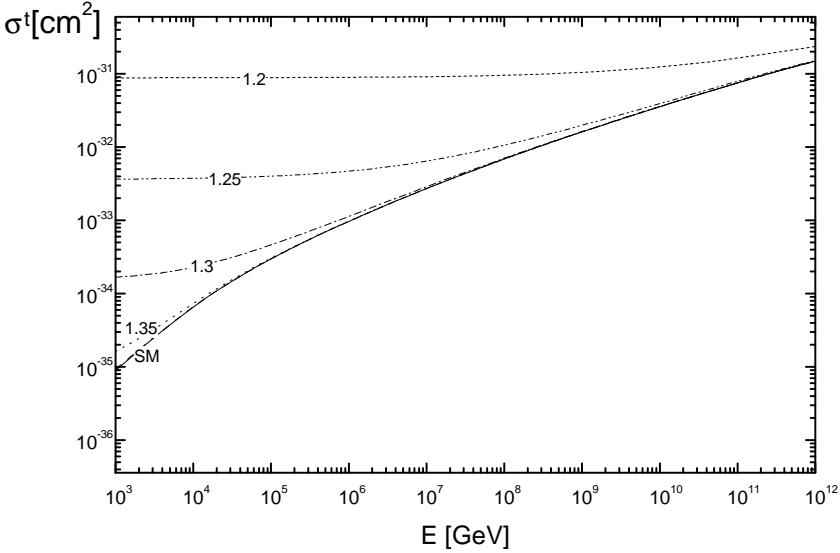


FIG. 2: Total cross section for the SM and for different values of the unparticles dimension  $d_u$  and  $c_L = c_R = 0.01$ .

where the corresponding parton distributions for a isoscalar target read

$$\begin{aligned} Q^U(x, Q^2) &= \frac{u_v(x, Q^2) + d_v(x, Q^2)}{2} + \frac{u_s(x, Q^2) + d_s(x, Q^2)}{2} \\ &\quad + c_s(x, Q^2) + t_s(x, Q^2) \\ Q^D(x, Q^2) &= \frac{u_v(x, Q^2) + d_v(x, Q^2)}{2} + \frac{u_s(x, Q^2) + d_s(x, Q^2)}{2} \\ &\quad + s_s(x, Q^2) + b_s(x, Q^2) \\ \bar{Q}^U(x, Q^2) &= \frac{u_s(x, Q^2) + d_s(x, Q^2)}{2} + c_s(x, Q^2) + t_s(x, Q^2) \\ \bar{Q}^D(x, Q^2) &= \frac{u_s(x, Q^2) + d_s(x, Q^2)}{2} + s_s(x, Q^2) + b_s(x, Q^2). \end{aligned}$$

In Fig. 2 we show the behavior of the total cross section ( $\sigma^t(E) = \sigma^{CC}(E) + \sigma^{NC}(E)$ ) with the neutrino energy for different values of  $d_u$  and  $c_L = c_R = 0.01$ . We can appreciate a considerable disagreement with the SM predictions, due to the unparticle propagator, particularly for low values of  $d_u$  and low neutrino energy. This very disparate behavior do not directly translate to the neutrino flux due to strong regeneration effects, as we will see in the next section.

### III. THE SURVIVING NEUTRINO FLUX

The surviving flux of neutrinos of energy  $E$ , with inclination  $\theta$  with respect to nadir direction,  $\Phi(E, \theta)$ , is the solution of the complete transport equation [8]:

$$\frac{d \ln \Phi(E, \tau')}{d \tau'} = -\sigma^t(E) + \int_E^\infty dE' \frac{\Phi(E', \tau')}{\Phi(E, \tau')} \frac{d\sigma^{NC}}{dE}, \quad (16)$$

where the first term correspond to absorption effects and the second one to the regeneration. Here,  $\tau = \tau(\theta)$  is the number of nucleons per unit area in the neutrino path through the Earth,

$$\tau(\theta) = N_A \int_0^{2R_E \cos \theta} \rho(z) dz, \quad (17)$$

In order to find a solution for this equation we make the following approximation [9]: we replace the fluxes ratio inside the integral of the second member by the ratio of fluxes that solve the homogeneous equation (i.e., only considering absorption effects)

$$\frac{\Phi(E', \tau')}{\Phi(E, \tau')} \rightarrow \frac{\Phi_0(E')}{\Phi_0(E)} e^{-\Delta(E', E)\tau'} \quad (18)$$

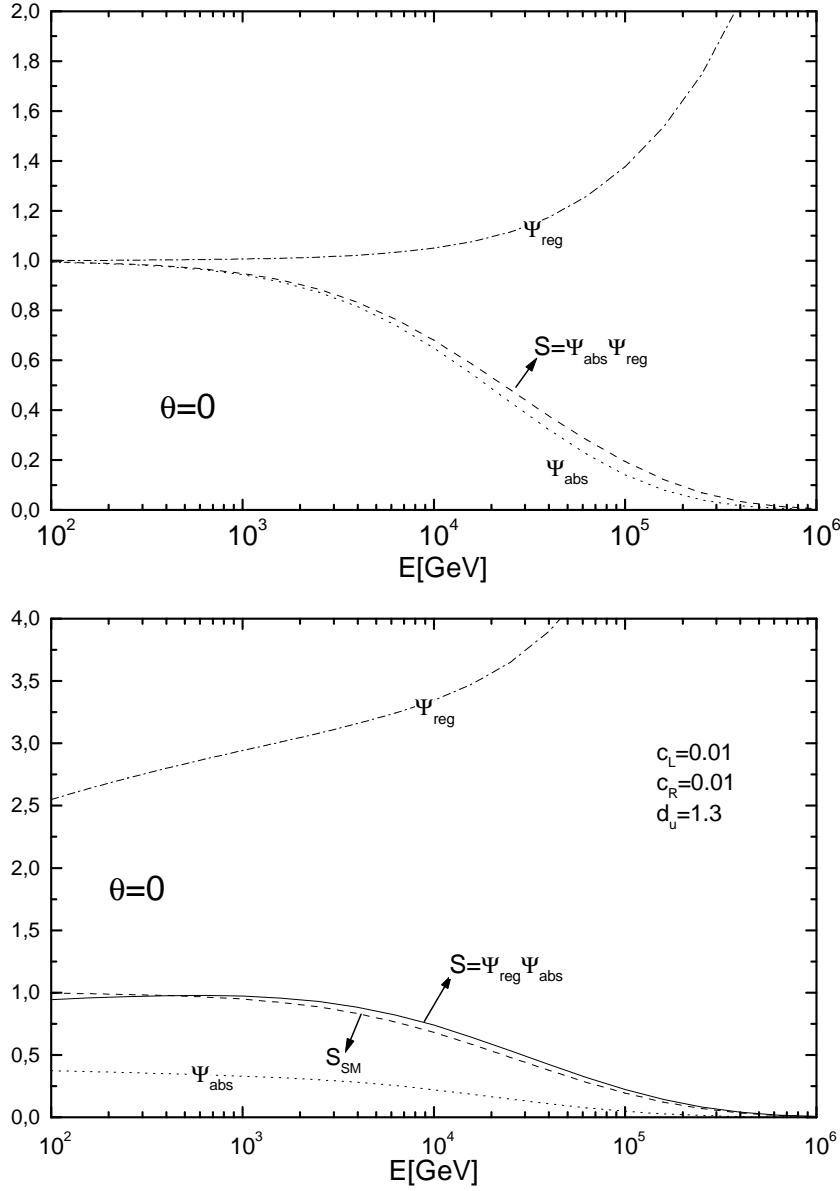


FIG. 3: Upper:  $S$  in the SM. We also include the factors  $\Psi_{\text{abs}}$  and  $\Psi_{\text{reg}}$ . Lower:  $S$  including the unparticle contribution for  $d_u = 1.3$  and  $c_L = c_R = 0.01$ .

where

$$\Delta(E', E) = [\sigma^t(E') - \sigma^t(E)] \tau' \quad (19)$$

Thus, integrating the transport equation we have

$$\Phi(E, \theta) = \Phi_0(E) e^{-\sigma_{\text{eff}}(E, \tau(\theta)) \tau(\theta)}, \quad (20)$$

where

$$\sigma_{\text{eff}}(E, \tau) = \sigma^t(E) - \sigma^{reg}(E, \theta), \quad (21)$$

with

$$\sigma^{reg}(E, \theta) = \int_E^\infty dE' \frac{d\sigma^{NC}}{dE'} \left( \frac{\Phi_0(E')}{\Phi_0(E)} \right) \left( \frac{1 - e^{-\Delta(E', E)\tau}}{\tau \Delta(E', E)} \right). \quad (22)$$

$\Phi_0(E)$  is the initial neutrino flux considered as isotropic,  $N_A$  is the Avogadro number,  $R_E$  is the radius of the Earth,  $\theta$  is the nadir angle taken from the down-going normal to the neutrino telescope and the earth density  $\rho(r)$  is given by the preliminary reference earth model [13]. It is important to mention that the solution of the transport equation, Eq. (20) is the first, but quite accurate, approximation of the iterative method showed in Ref.[10].

In order to illustrate the general behavior of the solution we show in Fig. 3, for  $\theta = 0$ , the factor  $S$

$$S = \frac{\Phi(E, \theta)}{\Phi_0(E)} = \Psi_{abs}(E, \theta)\Psi_{reg}(E, \theta), \quad (23)$$

with

$$\begin{aligned} \Psi_{abs}(E, \theta) &= e^{-\sigma_t(E)\tau(\theta)} \\ \Psi_{reg}(E, \theta) &= e^{\sigma_{reg}(E, \theta)\tau(\theta)} \end{aligned} \quad (24)$$

In this figure we show  $S$  for the SM and for the unparticles contribution with  $d_u = 1.3$  and  $c_L = c_R = 0.01$  and we have explicitly included both factors ( $\Psi_{abs}$  and  $\Psi_{reg}$ ) to see the compensation between absorption and regeneration effects.

For  $d_u \geq 1.3$  the absorption and the regeneration practically compensate each other and, then, the values of  $S$  are very near to the SM values (Fig. 3 Upper). For  $d_u < 1.3$  the absorption and regeneration effects do not cancel that efficiently and then we have values of  $S$  that clearly differ from the SM value. In Fig. 4 we show  $S$  for different values of the unparticle parameters and for two angles between the directions of the neutrino beam and the nadir.

#### IV. THE OBSERVABLE $\alpha(E)$

The angle  $\alpha(E)$  and the related ratio  $\eta(E)$  introduced in Ref.[11] are the observable that we shall use in this paper in order to study the impact of the unparticle physics on neutrino detection in a neutrino telescope such as IceCube. By definition  $\alpha(E)$  is the angle that divides the Earth into two homo-event sectors. When neutrinos traverse the planet in their journey to the detector, they find different matter densities, and then, different number of nucleons to interact with. In this conditions, the number of neutrinos that finally arrive to the detector depends on the arrival directions, indicated by the angle  $\theta$  with respect to the nadir direction. If we consider only upward-going neutrinos of a given energy  $E$ , that is, the ones with arrival directions  $\theta$  such that  $0 < \theta < \pi/2$ , there will always exist an angle  $\alpha(E)$  such that the number of events for  $0 < \theta < \alpha(E)$  equals that for  $\alpha(E) < \theta < \pi/2$ .

Clearly, the value of  $\alpha(E)$  is energy dependent. For low energies, the cross section decreases and the Earth becomes transparent to neutrinos. In this case  $\alpha(E) \rightarrow \pi/3$  for a diffuse isotropic flux since this angle divides the hemisphere into two sectors with the same solid angle. Obviously for extremely high energies, where most neutrinos are absorbed,  $\alpha(E) \rightarrow \pi/2$ , and for intermediate energies  $\alpha(E)$  varies accordingly between these limiting behaviors.

In order to define  $\alpha(E)$  we consider the expected number of events (muon tracks though charged currents  $\nu_\mu N$  interactions) at IceCube in the energy interval  $\Delta E$  and in the angular interval  $\Delta\theta$  that can be estimated as

$$\mathcal{N} = n_T T \int_{\Delta\theta} \int_{\Delta E} d\Omega dE_\nu \sigma^{CC}(E) \Phi(E, \theta), \quad (25)$$

where  $n_T$  is the number of target nucleons in the effective detection volume,  $T$  is the running time, and  $\sigma^{CC}(E)$  is the charged neutrino-nucleon cross section. We take the detection volume for the events equal to the instrumented volume for IceCube, which is roughly  $1 \text{ km}^3$  and corresponds to  $n_T \simeq 6 \times 10^{38}$ . The function  $\Phi(E, \theta)$  in Eq.(25) is the survival flux which is the solution Eq.(20) of the complete transport equation [8].

The definition of  $\alpha(E)$  is essentially the equality between two number of events, thus, to a good approximation, for each energy bin all the previous factors cancel except the integrated fluxes at each side. In this way,  $\alpha(E)$  can be defined by the equation

$$\int_0^{\alpha(E)} d\theta \sin \theta e^{-\sigma_{eff}(E)\tau(\theta)} = \int_{\alpha(E)}^{\pi/2} d\theta \sin \theta e^{-\sigma_{eff}(E)\tau(\theta)}, \quad (26)$$

which is numerically solved to give the results shown in the Fig. 5. There we show the SM prediction for  $\alpha(E)$  and the unparticles contribution for different values of the dimension  $d_u$ .

The main characteristics of  $\alpha(E)$  have been reported recently in Ref.[11]. It is worth to notice that  $\alpha(E)$  is weakly dependent on the initial flux but, at the same time it is strongly dependent on the neutrino nucleon cross-section. Hence, the use of the observable  $\alpha(E)$  reduces the effects of the experimental systematics and initial flux dependence.

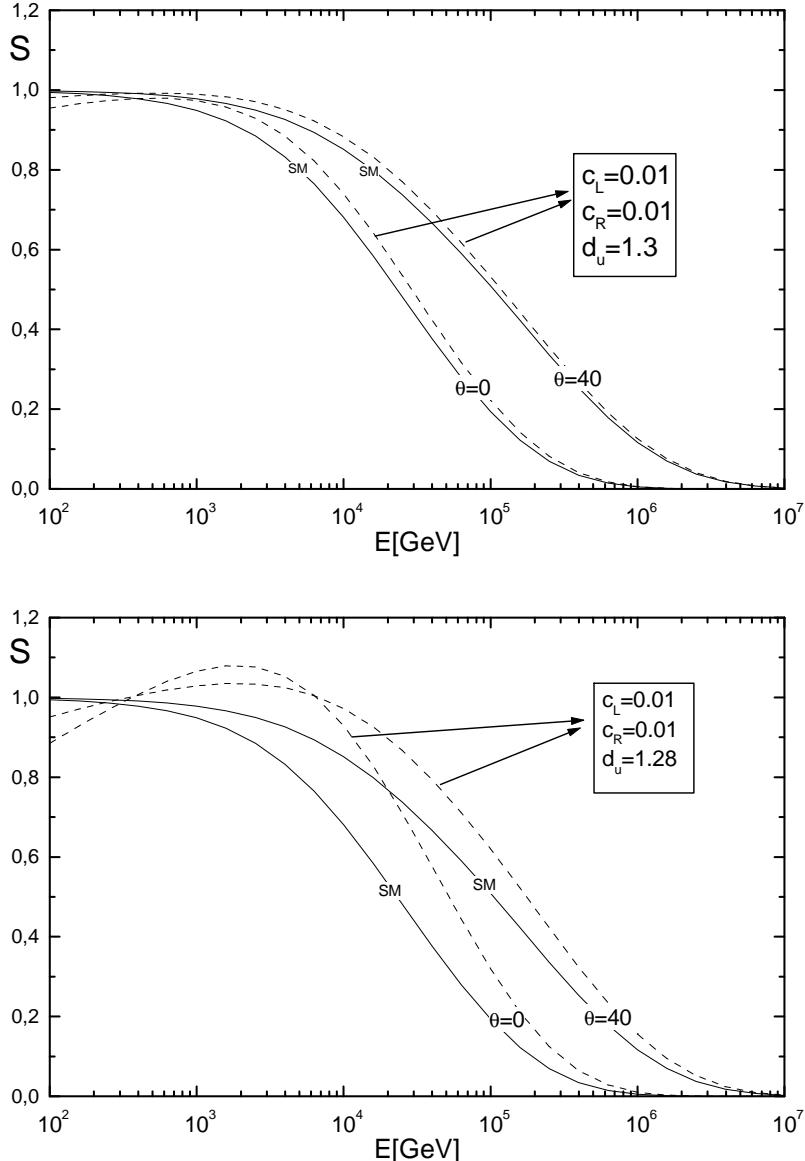


FIG. 4:  $S$  for the SM and for different values of  $d_u$  at two different angles with respect to nadir.

Since the functional form of  $\alpha(E)$  sharply depends on the interaction cross section neutrino-nucleon, if physics beyond the SM operates at these high energies, it will become manifest directly onto the function  $\alpha(E)$ .

In order to evaluate the impact of the observable  $\alpha(E)$  to bound new physics effects, we have estimated the corresponding uncertainties on the SM prediction for  $\alpha(E)$ . Considering the number of events as distributed according to a Poisson distribution the uncertainty can be propagated onto the angle  $\alpha_{\text{SM}}(E)$ . The number of events  $N$  as a function of  $\alpha_{\text{SM}}$  is

$$N = 2\pi n_T T \Delta E \sigma^{CC}(E) \Phi_0(E) \int_0^{\alpha_{\text{SM}}} d\theta \sin \theta e^{-\sigma_{\text{eff}}(E)\tau(\theta)}, \quad (27)$$

where we have considered the effective volume for contained events so that an accurate and simultaneous determination of the muon energy and shower energy is possible. For IceCube, it corresponds to the instrumented volume, roughly  $1 \text{ km}^3$ , implying a number of target nucleons  $n_T \simeq 6 \times 10^{38}$ . We have considered an integration time  $T = 15 \text{ yr}$  which is the expected lifetime of the experiment. To propagate the error on  $N$  to obtain the one on  $\alpha$ , we note that

$$\delta N = \frac{dN}{d\alpha} \delta \alpha, \quad (28)$$

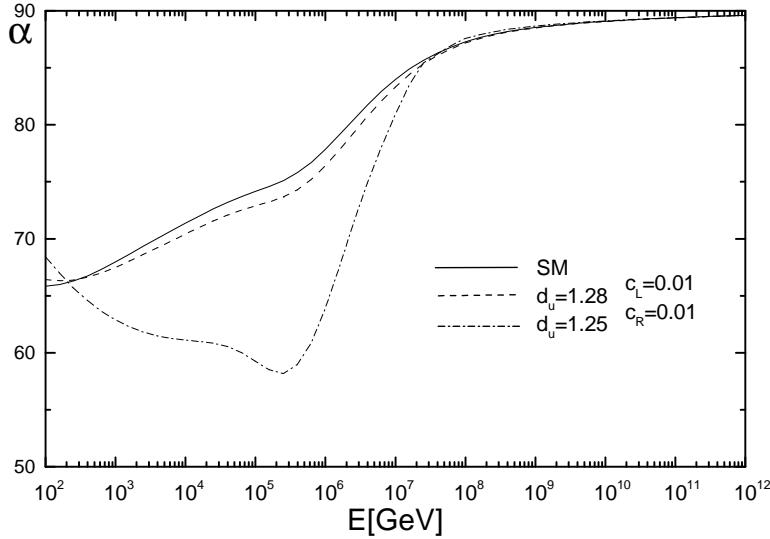


FIG. 5: The predictions for  $\alpha(E)$  obtained for different values on  $d_u$ .

and dividing by  $N$  we obtain

$$\delta\alpha = \left[ \int_0^{\alpha_{SM}(E)} d\theta \left( \frac{\sin\theta}{\sin\alpha_{SM}(E)} \right) e^{\sigma_T(E)[\tau(\alpha_{SM}(E)) - \tau(\theta)]} \right] \left( \frac{\delta N}{N} \right), \quad (29)$$

where for Poisson distributed events we have

$$\delta N = \sqrt{N}. \quad (30)$$

In order to evaluate the errors on  $\alpha(E)$ , it is necessary to consider a level of initial flux  $\Phi_0(E)$ . Here we have added together the cosmological diffuse isotropic flux and the atmospheric one (see Fig. 7). For the atmospheric flux, we have considered the one given in Ref.[14]. As for the cosmological diffuse flux, the usual benchmark is the so-called Waxman-Bahcall (WB) flux for each flavor,  $E_{\nu_\mu}^2 \phi_{WB}^{\nu_\mu} \simeq 2.4 \times 10^{-8} \text{ GeV cm}^{-2} \text{s}^{-1} \text{sr}^{-1}$ , which is derived assuming that neutrinos come from transparent cosmic ray sources [15], and that there is an adequate transfer of energy to pions following  $pp$  collisions. However, one should keep in mind that if there are in fact hidden sources which are opaque to ultra-high energy cosmic rays, then the expected neutrino flux will be higher.

On the other hand, we have the experimental bounds set by AMANDA. A summary of these bounds can be found in Refs.[16, 17] and as a representative value we take  $E_{\nu_\mu}^2 \phi_{AM}^{\nu_\mu} \simeq 8 \times 10^{-8} \text{ GeV cm}^{-2} \text{s}^{-1} \text{sr}^{-1}$ . With the intention to estimate the number of events, we have considered an intermediate flux (INT) level slightly below the present experimental bound by AMANDA,

$$E_{\nu_\mu}^2 \phi_{INT}^{\nu_\mu} \simeq 5 \times 10^{-8} \text{ GeV cm}^{-2} \text{s}^{-1} \text{sr}^{-1}. \quad (31)$$

As it was discussed in Ref.[11] the interval for maximum sensitivity for  $\alpha$  is  $10^5 \text{ GeV} < E < 10^7 \text{ GeV}$ . However, as for lower energies the atmospheric flux grows and then the errors fall, we have considered as an energy window for the fits the interval:  $10^3 \text{ GeV} < E < 10^7 \text{ GeV}$ . In Fig. 6 we show our results for the observable  $\alpha$  and the corresponding errors within the mentioned energy window. In Fig. 7 we show the used flux.

In the same context, we can define another observable related to  $\alpha(E)$ . We consider the hemisphere  $0 < \theta < \pi/2$  divided into two regions by the angle  $\alpha_{SM}(E)$ ,  $\mathcal{R}_1$  for  $0 < \theta < \alpha_{SM}(E)$  and  $\mathcal{R}_2$  for  $\alpha_{SM}(E) < \theta < \pi/2$ . We then calculate the ratio  $\eta(E)$  between the number of events for each region,

$$\eta(E) = \frac{N_1}{N_2}, \quad (32)$$

where  $N_1$  is the number of events in the region  $\mathcal{R}_1$  and  $N_2$  the number of events in the region  $\mathcal{R}_2$ . By using  $\eta(E)$  the effects of experimental systematic and initial flux dependence are reduced. If there is only SM physics, then we have that the ratio  $\eta_{SM}(E) = 1$ . In order to estimate the capability of  $\eta(E)$  to bound unparticle effects, we have considered the values of  $\eta(E)$  along with their error bars in Fig. 8 as if they had been obtained from experimental measurements for  $\eta(E)$ . We proceed, then, to perform a  $\chi^2$ -analysis taking as free parameters the dimension  $d_u$  and the constant  $c = c_L = c_R$  and considering as *experimental point* the SM values for  $\eta(E)$  for the same energy bin used in Fig. 6. We define the  $\chi^2$  function in the usual way,

$$\chi^2 = \sum_{i=1,8} \frac{(\eta_{SM}(E_i) - \eta(E_i, d_u, c))^2}{(\delta\eta(E_i))^2}. \quad (33)$$

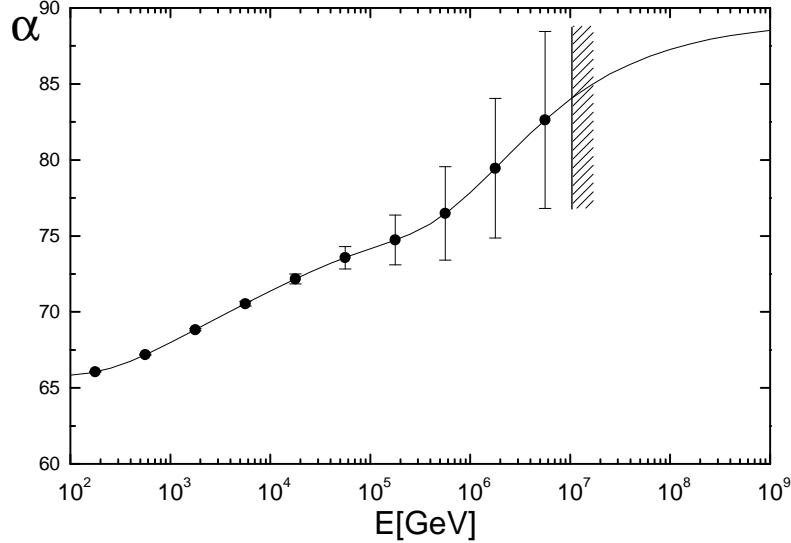


FIG. 6: The prediction for  $\alpha_{SM}(E)$  and the statistical errors.

where according to the definition of  $\eta(E)$  (Eq.(32)) the statistical errors are given by  $\delta\chi(E_i) = 2/\sqrt{N_i}$  for events distributed according to a Poisson distribution. The function  $\chi^2$  is minimized to obtain the allowed region in the  $(d_u, c)$  plane, which corresponds to the region below the curve **A** shown in Fig. 9. In the same figure we also include other bounds obtained from different processes. The curve **B** corresponds to bounds obtained from atomic parity violation[18], the curve **C** corresponds to bounds from the muon anomaly [19] and curve **E** comes from low energy  $\nu_e - e$  scattering[20]. This last bound is significatively better than all the other bounds. We would like to stress that the curve **E** corresponds to  $\nu_e - e$  scattering but we are considering muon neutrinos, i.e., a different neutrino flavor. Thus, if we consider the unparticle interactions as non-universal then this constraint would not apply and the neutrino telescope bounds may be competitive. However, there are experimental data for the ratio between  $\nu_\mu - N$  Neutral Current to Charged Current interaction  $R$  [21]:

$$R(E_{\nu_\mu}) = \frac{\sigma_{\nu N \rightarrow \nu X}^{NC}(E)}{\sigma_{\nu N \rightarrow \mu X}^{CC}(E)} \quad (34)$$

In order to compare with the experimental value

$$R = 0.320 \pm 0.010, \quad (35)$$

which is an average measurement for energies in the range 20 GeV - 200 GeV, we calculated the corresponding average value for  $R$ :

$$\langle R \rangle = \frac{1}{180 \text{GeV}} \int_{20 \text{GeV}}^{200 \text{GeV}} R(E) dE \quad (36)$$

and compared it with the previous experimental value (Eq.(35)) for different values of the unparticles parameters  $(d_u, c)$ . The allowed region corresponds to the points below the curve **D** of Fig. 9. This curve shows that the bound coming from  $R$  is stronger than the one IceCube could possible set in the future.

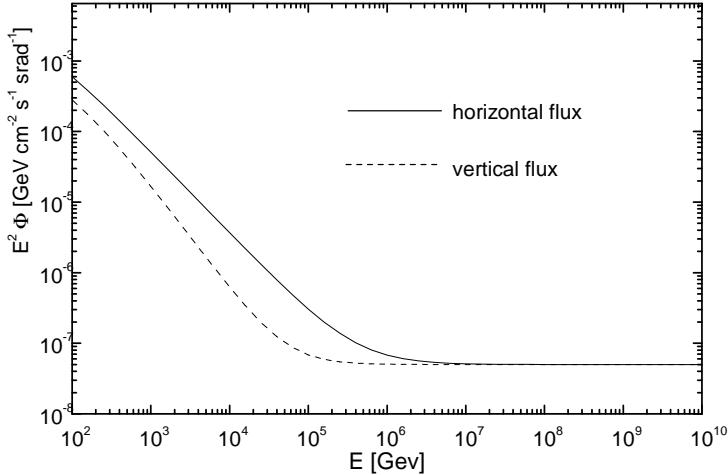


FIG. 7: The utilized flux obtained by adding the atmospheric and the isotropical cosmic flux.

## V. CONCLUSIONS

In the present work we studied the effects of unparticles contributions to the neutrino-nucleon cross section on the survival neutrino flux in a neutrino telescope like IceCube. To do it, we considered an effective interaction between standard particles and unparticles. We have found a considerable disagreement with the SM prediction for the neutrino observables defined above, particularly for low values of  $d_u$  and low neutrino energy. For moderate values of  $d_u$  this disagreement tends to disappear due to a strong cancellation between absorption and regeneration.

We have also studied the possibility to bound effects of unparticles contributions to the interactions between muon neutrinos and the nucleons of the Earth using the observable  $\eta(E)$ . In this context, we fitted the theoretical expression for  $\eta(E)$  as a function of the  $d_u$  and  $c$  taking as experimental data the SM values obtained for  $\eta$  ( $\eta_{SM}(E) = 1$ ) along with the errors derived for a number of events distributed according to a Poisson distribution. The results are shown in Fig. 9 as a allowed region plot. Finally, in the same figure we have compared the limits that have been obtained for different authors for low energy processes. The bounds coming from the ratio R of neutral to charged currents (curve D) is more restrictive than the bound IceCube could set (curve A) and then new effects that IceCube may detect could not be attributable to unparticle physics.

## Acknowledgments

We thank CONICET (Argentina) and Universidad Nacional de Mar del Plata (Argentina), the Fundación del Banco de la Republica (Colombia), PDT 54/094 (Uruguay) and CLAF for the financial supports.

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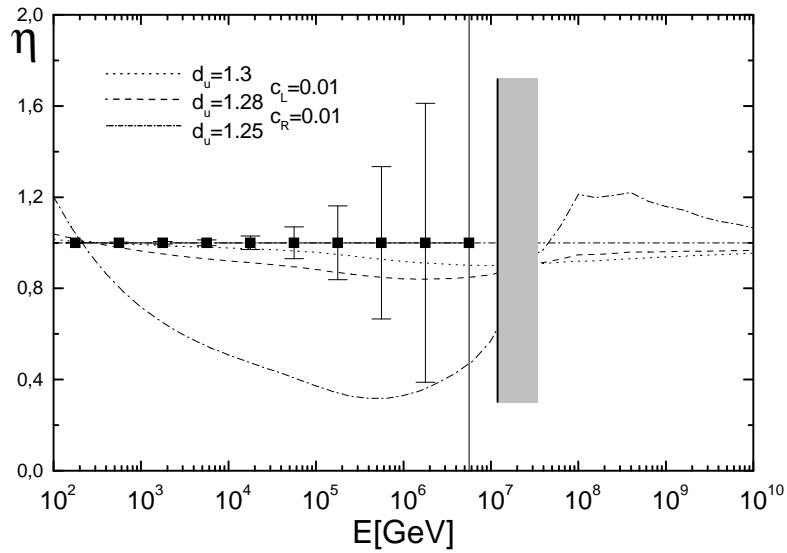


FIG. 8:  $\eta(E)$  for different values of  $d_u$ . We include the statistical errors obtained of a number of events distributed as a Poisson distribution.

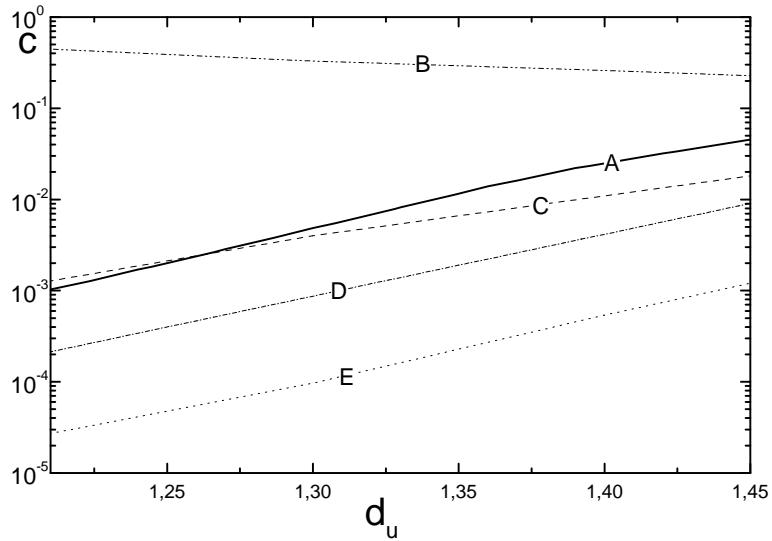


FIG. 9: The allowed regions in the  $(d_u, c)$  plane are below the corresponding curves. Curve A: this paper, Curve B: bounds obtained form Atomic Parity Viloation [18], Curve C: it obtained from  $(g - 2)_\mu$  [19], Curve E correspond to bounds obtained from  $\nu_e - e$  scattering [20], and curve **D** correspond to ones obtained from the ratio  $\langle R \rangle$  [21].

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